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The Photon Wave Function in Non-forward Diffractive Scattering with Non-vanishing Quark Masses

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Abstract

We present the Photon wave function for explicit helicity states, valid for massive quarks, both in momentum and configuration space. We further investigate diffractive scattering at nonzero momentum transfer and find, that we have a similar factorisation as in the case of zero momentum transfer.

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1.) In this letter we want to consider the amplitude for the diffractive process $\gamma^* p \to p + X$ both in momentum and configuration space. We consider the kinematic limit, where the $\gamma^* p$ energy s is much larger than the diffractive mass M_X^2, the relative transverse momentum squared of the $q\bar{q}$-pair $\mathbf{k}^2$, the virtuality of the photon Q^2 and the momentum transfer t from the virtual photon system to the proton. In this limit the amplitude is factorised basically into the photon wave function and the unintegrated off-diagonal gluon structure function.

Instead of the space of transverse momenta one can transform this amplitude to configuration space, where the conjugate variable to the transverse momentum separation of the $q\bar{q}$ pair is its transverse size. This leads to the convenient 'dipole-picture' in the proton rest frame, saying that the virtual photon fluctuates into the $q\bar{q}$-pair, the colour dipole, long before it undergoes the interaction with the proton and therefore this (perturbative) fluctuation and the interaction of the dipole with the proton are independent and factorise in the transverse space. We are left with the photon wave function, describing the independent fluctuation into colour dipoles, and the cross section for the dipole scattering off the proton. In this way a description of cross sections for inclusive and diffractive scattering processes as well as of diffractive vector meson production can be carried out completely in terms of the photon wave function and the so called dipole cross section, see e.g. [1, 2, 3].

The concept of the photon wave function is rather well known. In [4] expressions for massless quarks and transverse photons are given explicitly in terms of helicity amplitudes. [5, 6] give explicit expressions for the squared wave functions, including longitudinal photons and also massive quarks. A derivation of the amplitudes in terms of an explicit spinor representation limited to massless quarks is given in [7].

We will therefore reconsider a perturbative derivation of the wave function, including massive fermions. This approach is then suitable to give a rather general expression for the diffractive amplitude, i.e. we do not take the limit of vanishing momentum transfer t, as usual. In momentum space, this expression still exhibits similar properties as the expression for the diffractive amplitude at $t \to 0$. Going to configuration space, aside of the transverse separation of the colour dipole, we have to introduce the impact parameter, giving the transverse distance between the incoming dipole and the proton. Also here the picture, that the fluctuation of the virtual photon into the dipole and the interaction of the dipole with the proton are independent, will hold.

2.) Before considering the amplitudes for diffractive scattering, where the virtual photon couples, after fluctuation into a $q\bar{q}$-pair, to two t-channel gluons, projected into colour singlet, we calculate the photon wave function from the simpler process, where a single gluon is exchanged in the t-channel as depicted in Fig. 1.

Throughout the calculation we will use a Sudakov decomposition of all momenta with respect to the light cone vectors $q' = q - (Q^2/s)p_B$ and $p' = p_B + (m_B^2/s)q'$; $s = 2p' \cdot q'$; q is the momentum of the photon with virtuality $Q^2 = -q^2$ and p_B is the momentum of the incoming quark B with mass m_B. For the internal quark line in diagram (a) we use the momentum k and for the exchanged gluon the momentum r, heading towards the quark-antiquark system,

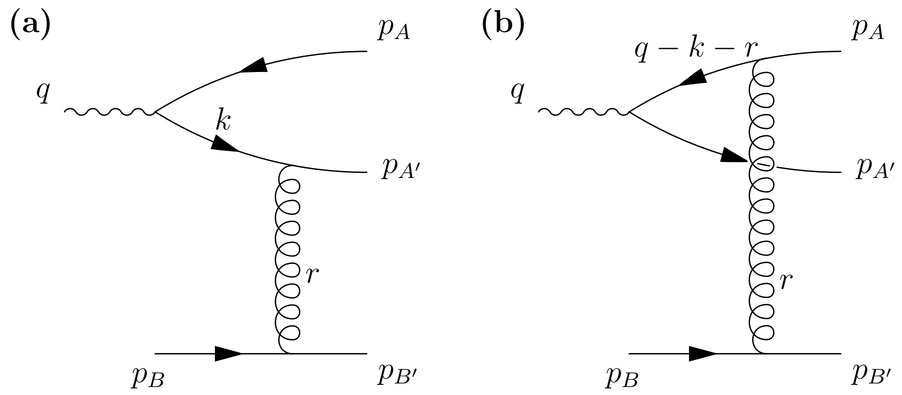

Figure 1: The Feynman diagrams contributing to $\gamma^* + q' \longrightarrow q\bar{q} + q'$ in the high energy limit.

having the following Sudakov decomposition:

$$k = \alpha q' + \beta p' + k_\perp, \tag{1}$$
$$r = \frac{t}{s} q' + x_{\mathbb{P}} p' + r_\perp. \tag{2}$$

In this way the antiquark and quark carry fractions α and $(1-\alpha)$ in the q'-direction. In the following we often use the Euclidean part of the transverse momenta explicitly, e.g. $k_\perp^2 = -\mathbf{k}^2$. The transverse momenta $-\mathbf{k}$ and $\mathbf{k}+\mathbf{r}$ are carried by the antiquark and quark, respectively, against the proton that recoils with $-\mathbf{r}$. From the on-shell conditions for the outgoing quarks, together with the assumptions that $Q^2 \sim |r^2 = t| \sim M_X^2 \ll s$, we find $x_{\mathbb{P}}$ and β to be of the order t/s.

In this limit we use the usual decomposition of the metric tensor

$$g_{\mu\nu} = \frac{2}{s}(p'_\mu q'_\nu + p'_\nu q'_\mu) + g^\perp_{\mu\nu} \tag{3}$$

for the gluon propagator and retain only the first term with an appropriate contraction, while the other terms are suppressed by powers of t/s. Furthermore, We express the denominators of the quark lines in diagrams (a) and (b) respectively as

$$\Delta_a = k^2 - m^2 = -\frac{1}{1-\alpha}[\alpha(1-\alpha)Q^2 + \mathbf{k}^2 + m^2] \equiv -\frac{1}{1-\alpha} D(\mathbf{k}), \tag{4}$$
$$\Delta_b = (q-k-r)^2 - m^2 = -\frac{1}{\alpha}[\alpha(1-\alpha)Q^2 + (\mathbf{k}+\mathbf{r})^2 + m^2] \equiv -\frac{1}{\alpha} D(\mathbf{k}+\mathbf{r}).$$

The mass in the upper fermion line is always denoted as m.

Before writing down the amplitude we first simplify the contribution from the lower quark line, which is well known to give

$$\bar{u}(p_{B'})\gamma^\mu u(p_B) \approx (p_B + p_{B'})^\mu \delta_{\lambda_B, \lambda_{B'}}. \tag{5}$$

contracting with q' from the gluon propagator, we basically get $s\delta_{\lambda_B,\lambda_{B'}}$. A helicity flip term is suppressed by a power of $m_B/\sqrt{s}$, as we have also checked with the helicity method from [8], that we used in the remainder of the calculation.

After this we may write down the amplitude for one gluon exchange as

$$\mathcal{A} = -ee_f\mathcal{C}\frac{2g^2\delta_{\lambda_B,\lambda_{B'}}}{t}\bar{u}_{\lambda'}(p_{A'})\left(\frac{\chi_a}{\Delta_a}+\frac{\chi_b}{\Delta_b}\right)v_\lambda(p_A) \tag{6}$$

with

$$\chi_a = \not{p}'(\not{k}+m)\not{\varepsilon}^\gamma, \qquad \chi_b = -\not{\varepsilon}^\gamma(\not{q}-\not{k}-\not{r}-m)\not{p}', \tag{7}$$

arising from the upper fermion lines and $\lambda'(\lambda) = \pm$ denotes the helicity of the (anti-)quark. Within these we choose the polarisation vector of the virtual photon as

$$\varepsilon^0 = \frac{1}{Q}(q'+x_Bp'), \qquad \varepsilon_\perp^\gamma = \frac{1}{\sqrt{2}}(0,1,\gamma i,0). \tag{8}$$

Here and in the following $\gamma = 0,\pm$ always denotes the photon polarisation. In eq. (6) ee_f gives the quark charge, g is the strong coupling constant and $\mathcal{C}$ is the colour factor.

What remains to be done is to evaluate helicity amplitudes of the general form $\bar{u}_{\lambda'}(p_{A'})\chi v_\lambda(p_A)$. Since we want to allow for quarks of non-vanishing mass m here, the use of an explicit spinor representation might be much more complicated as in the case of massless fermions. Therefore, we make use of the helicity method given in [8]. The basic idea of this method is to express spinors for massive particles in terms of massless spinors and to observe, that the above helicity amplitude can be written as a trace:

$$\bar{u}_{\lambda'}(p_{A'})\chi v_\lambda(p_A) = \mathrm{tr}\,\chi v_\lambda(p_A)\bar{u}_{\lambda'}(p_{A'}). \tag{9}$$

With the help of massless spinors it is then possible to express the matrix $v_\lambda(p_A)\bar{u}_{\lambda'}(p_{A'})$ explicitly in terms of the fermions' helicities and momenta and two reference vectors k_0, k_1, of which the final result will not depend. After several steps the expressions for longitudinal and transverse polarisations of the photon for diagram (a), and a similar one for diagram (b), are obtained,

$$\chi_a(\gamma=0) = \alpha s\,\not{\varepsilon}^0 + \frac{s}{2Q}\,\not{r}_\perp - i\epsilon(\mu,\varepsilon^0,p_{B'},p_B)\gamma_\mu\gamma_5, \tag{10}$$

$$\chi_a(\gamma=\pm) = \alpha s\,\not{\varepsilon}_\perp + \boldsymbol{\varepsilon}_\perp\cdot\mathbf{r}\,\not{p}' - i\epsilon(\mu,\varepsilon_\perp,p_{B'},p_B)\gamma_\mu\gamma_5. \tag{11}$$

Here we only kept terms that will be present in the high energy limit. Applying the helicity method we finally obtain for the amplitude (6) a result, which can be written in a compact form

$$\mathcal{A} = eg^2\frac{2s}{t}\delta_{\lambda_B,\lambda_{B'}}\mathcal{C}\sqrt{\alpha(1-\alpha)}\Big(\Psi(\mathbf{k},\alpha)-\Psi(\mathbf{k}+\mathbf{r},\alpha)\Big). \tag{12}$$

Here, $\Psi(\mathbf{k},\alpha)$ is the so called photon wave function, which contains all dependence of the amplitude on the properties of the virtual photon and the $q\bar{q}$-pair, including helicities. With the notation $\Psi^{\gamma}_{\lambda'\lambda}(\mathbf{k})$ for the explicit helicity dependence of the wave function and $\underline{k} = k_x + i\gamma k_y$ for the transverse momenta, we get the following explicit results:

$$\Psi^0_{\pm\mp}(\mathbf{k},\alpha) = \frac{2e_f\alpha(1-\alpha)Q}{\alpha(1-\alpha)Q^2+\mathbf{k}^2+m^2}, \tag{13}$$

$$\Psi^{\pm}_{\pm\mp}(\mathbf{k},\alpha) = \frac{\sqrt{2}e_f\alpha\underline{k}}{\alpha(1-\alpha)Q^2+\mathbf{k}^2+m^2}, \tag{14}$$

$$\Psi^{\pm}_{\mp\pm}(\mathbf{k},\alpha) = \frac{-\sqrt{2}e_f(1-\alpha)\underline{k}}{\alpha(1-\alpha)Q^2+\mathbf{k}^2+m^2}, \tag{15}$$

$$\Psi^{\pm}_{\pm\pm}(\mathbf{k},\alpha) = \frac{\sqrt{2}e_f im}{\alpha(1-\alpha)Q^2+\mathbf{k}^2+m^2}, \tag{16}$$

$$\Psi^0_{\pm\pm}(\mathbf{k},\alpha) = \Psi^{\pm}_{\mp\mp}(\mathbf{k},\alpha) = 0. \tag{17}$$

In the additional term (16), appearing for massive quarks, we note a naive helicity conservation $\gamma = (\lambda+\lambda')/2$. A similar situation was found for (13) but here the mass dependence always cancels in the difference of the wave functions. On the other hand, in the massless limit we can only have the configuration $\lambda = -\lambda'$ within the $q\bar{q}$-pair. This behaviour is easily understood from the appropriate limit in $m/\sqrt{s}$: in the high energy limit the fermions' helicities are always opposite and independent of the coupling, in contrast to the nonrelativistic limit where we have helicity conservation.

Turning to diffractive scattering we can determine the photon part of a diffractive amplitude from the exchange of two gluons in the t-channel projected into colour singlet and coupling in all possible ways to the $q\bar{q}$-pair as depicted in Fig. 2. In the leading $\log x$ approximation the amplitude is dominated by its imaginary part, from which we will reconstruct the full amplitude later on. By virtue of the Cutkosky rules we may consider each diagram as a composition of two, where the crossed fermion lines shown in Fig. 2 are on mass shell. A left part, where we have a structure as in the one gluon exchange, discussed above, and a right part being just a simple one gluon exchange between two quark lines. From the cuts in the fermion lines we obtain similar kinematics for the loop momentum

$$\ell = \alpha_\ell q' + \beta_\ell p' + \ell_\perp \tag{18}$$

as for the momentum transfer r in the one gluon exchange, i.e. the longitudinal components α_ℓ and β_ℓ are of the order t/s and in the loop there is basically transverse momentum circulating. Therefore, in the left part of the diagrams we have the same kinematical conditions as in the one gluon exchange case.

From the respective right parts of the diagrams we get a factor $2\alpha s/(\mathbf{r}-\boldsymbol{\ell})^2$ or $2(1-\alpha)s/(\mathbf{r}-\boldsymbol{\ell})^2$ with helicity conservation within the fermion lines, depending on whether the right gluon couples to the quark or the antiquark. On the other hand we get factors $1/2\alpha s$ or $1/2(1-\alpha)s$ from the integration over the longitudinal part of the loop momentum and the on-shell conditions, therefore we are only left with the denominators from the right parts of the diagrams.

To determine the left parts, we can simply read off the transverse momentum of the virtual quark line from the diagrams and get a photon wave function with this argument, similar to

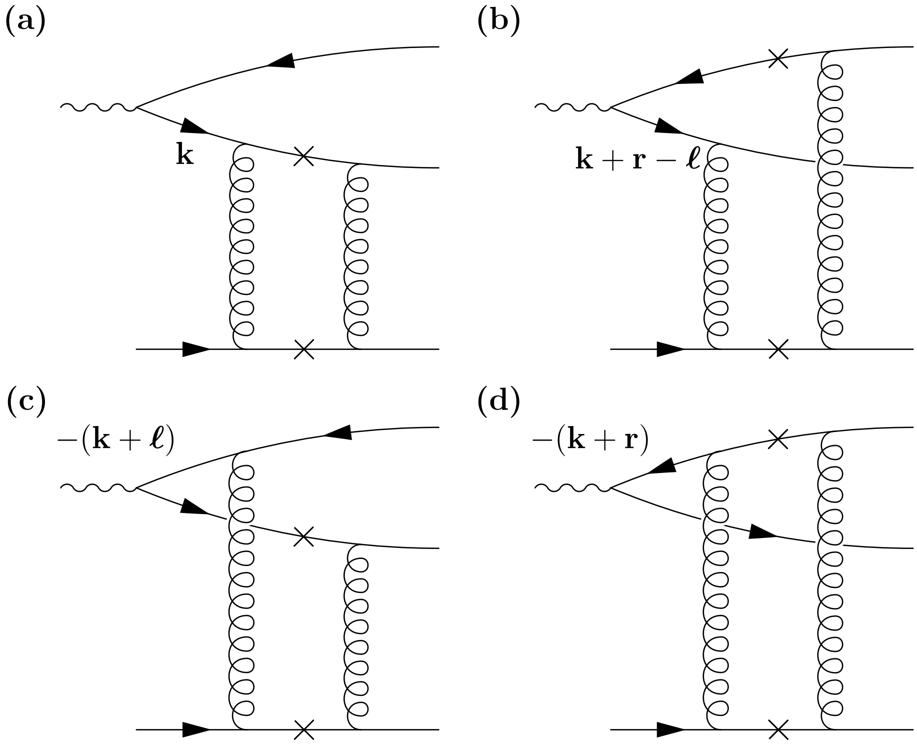

Figure 2: The Feynman diagrams contributing to diffractive $\gamma^* p \longrightarrow q\bar{q}p$ scattering. Only the transverse momenta of virtual fermion lines are denoted.

the one gluon case, but now we also have the transverse part of the loop momentum as an argument of the wave functions. Furthermore we now have a different colour structure, of which it is most important for us, that in the colour singlet case, we are interested in, the two gluons are symmetric in colour. Putting everything together, we can write the full amplitude, arising from the diagrams in Fig. 2 as

$$\begin{aligned}\mathcal{A} = ieg^4 \frac{s}{t} \delta_{\lambda_B,\lambda_{B'}} \mathcal{C}' \sqrt{\alpha(1-\alpha)} \int \frac{d^2\boldsymbol{\ell}}{(2\pi)^2} \frac{\mathbf{r}^2}{\boldsymbol{\ell}^2(\mathbf{r}-\boldsymbol{\ell})^2} \\ \times \Big\{ \Psi(\mathbf{k},\alpha) + \Psi(\mathbf{k}+\mathbf{r},\alpha) - \Psi(\mathbf{k}+\boldsymbol{\ell},\alpha) - \Psi(\mathbf{k}+\mathbf{r}-\boldsymbol{\ell},\alpha) \Big\},\end{aligned} \tag{19}$$

with the wave function $\Psi(\mathbf{k},\alpha)$ defined as in eqns. (13)–(17). In the limit $\mathbf{r}^2 = -t \to 0$ for the term in brackets we obtain the same expression of a 'double difference' as in [9].

To include the non-perturbative coupling of the two gluons to the proton, we have to use an unintegrated off-diagonal gluon distribution $\mathcal{F}(x, x', \boldsymbol{\ell}^2, \mathbf{r}^2)$, a suitable generalisation of the unintegrated gluon distribution to the off-diagonal case, see e.g. [10] for a discussion of off-diagonal gluon distributions in diffraction. Here, x and x' denote the longitudinal momentum fractions of the two gluons, coupling to the proton and $x_{\mathbb{P}} = x - x'$. Integrating over $\boldsymbol{\ell}$ would give an off-diagonal gluon distribution,

$$\int^{Q^2} d^2\boldsymbol{\ell}\,\mathcal{F}(x,x',\boldsymbol{\ell}^2,\mathbf{r}^2) = G(x,x',\mathbf{r}^2,Q^2) \tag{20}$$

and in the limit $x \approx x'$ we have $G(x,x,\mathbf{r}^2=0,Q^2) = xg(x,Q^2)$, with $g(x,Q^2)$ being a diagonal gluon distribution. In [11, 12] the ratio $G(x,x',t,Q^2)/xg(x,Q^2)$ is discussed in detail.

Now we may write the general amplitude for diffractive scattering off the proton, if we introduce the unintegrated off-diagonal gluon structure function $\mathcal{F}(x,x',\boldsymbol{\ell}^2,\mathbf{r}^2)$ as

$$\mathcal{A} = i\frac{\pi}{4}eg^2\sqrt{\alpha(1-\alpha)}s\int\frac{d^2\boldsymbol{\ell}}{\pi\boldsymbol{\ell}^2}\mathcal{F}(x,x',\boldsymbol{\ell}^2,\mathbf{r}^2)D\Psi(\mathbf{k},\mathbf{r},\boldsymbol{\ell}), \tag{21}$$

where we have introduced the shorthand notation

$$D\Psi(\mathbf{k},\mathbf{r},\boldsymbol{\ell}) = \Psi(\mathbf{k},\alpha) + \Psi(\mathbf{k}+\mathbf{r},\alpha) - \Psi(\mathbf{k}+\boldsymbol{\ell},\alpha) - \Psi(\mathbf{k}+\mathbf{r}-\boldsymbol{\ell},\alpha) \tag{22}$$

for the double difference of the wave functions. $\mathcal{F}(x,x',\boldsymbol{\ell}^2,\mathbf{r}^2)$ is normalised as in the diagonal case, i.e. $\mathcal{F}(x,x,\boldsymbol{\ell}^2,\mathbf{r}^2=0) = \mathcal{F}(x,\boldsymbol{\ell}^2)$, with $\mathcal{F}(x,\boldsymbol{\ell}^2)$ being the unintegrated gluon structure function used in [9, 13]. Furthermore, we left the dependence on t in the unintegrated structure function to keep our discussion general.

3.) It is more instructive to consider the factorisation of diffractive amplitudes in configuration space. Therefore we will Fourier-transform with respect to the transverse momenta $\mathbf{k}$ and $\mathbf{r}$. The variable conjugate to $\mathbf{k}$, the momentum separation of the $q\bar{q}$-pair, will be the transverse separation of the $q\bar{q}$-pair $\boldsymbol{\varrho}$, or simply the 'dipole size'. On the other hand, since we now include also a momentum transfer between the diffractive system and the proton, we have the impact parameter $\mathbf{b}$ between the virtual photon and the proton as the variable conjugated to the momentum transfer.

We get the conjugated photon wave functions $\psi^{\gamma}_{\lambda'\lambda}(\boldsymbol{\varrho})$ in $\boldsymbol{\varrho}$-space as

$$\psi^{0}_{\pm\mp}(\boldsymbol{\varrho}) = \frac{1}{\pi}e_f\alpha(1-\alpha)QK_0(\delta\varrho), \tag{23}$$

$$\psi^{\pm}_{\pm\mp}(\boldsymbol{\varrho}) = \frac{i}{\pi}e_f\alpha\delta\frac{\boldsymbol{\varrho}\cdot\boldsymbol{\varepsilon}}{\varrho}K_1(\delta\varrho), \tag{24}$$

$$\psi^{\pm}_{\mp\pm}(\boldsymbol{\varrho}) = -\frac{i}{\pi}e_f(1-\alpha)\delta\frac{\boldsymbol{\varrho}\cdot\boldsymbol{\varepsilon}}{\varrho}K_1(\delta\varrho), \tag{25}$$

$$\psi^{\pm}_{\pm\pm}(\boldsymbol{\varrho}) = \frac{im}{2\pi}e_fK_0(\delta\varrho), \tag{26}$$

where we used $\delta^2 = \alpha(1-\alpha)Q^2 + m^2$ as a shorthand notation and $K_\nu(z)$ is a modified Bessel function. After squaring, our results agree with that of [5]. Using these wave functions, we can easily transform the diffractive amplitude (21) into configuration space,

$$\tilde{\mathcal{A}}^D(\boldsymbol{\varrho},\mathbf{b}) = \int\frac{d^2\mathbf{k}}{(2\pi)^2}\int\frac{d^2\mathbf{r}}{(2\pi)^2}e^{i\mathbf{k}\cdot\boldsymbol{\varrho}}e^{i\mathbf{r}\cdot\mathbf{b}}\mathcal{A}^D(\mathbf{k},\mathbf{r}) \tag{27}$$

$$= B\int\frac{d^2\mathbf{k}}{(2\pi)^2}\int\frac{d^2\mathbf{r}}{(2\pi)^2}e^{i\mathbf{k}\cdot\boldsymbol{\varrho}}e^{i\mathbf{r}\cdot\mathbf{b}}\int\frac{d^2\boldsymbol{\ell}}{\pi\boldsymbol{\ell}^2}\alpha_s(\mu^2)\mathcal{F}(x,x',\boldsymbol{\ell}^2,\mathbf{r}^2)D\Psi(\mathbf{k},\mathbf{r},\boldsymbol{\ell}), \tag{28}$$

where B contains the factors in front of the integral in (21). In the $\mathbf{k}$ integration we make use of the wave function (23)-(26) and pick up appropriate phases from a shift in the integration region and get

$$\int \frac{d^2\mathbf{k}}{(2\pi)^2} e^{i\mathbf{k}\cdot\boldsymbol{\varrho}} D\Psi(\mathbf{k},\mathbf{r},\boldsymbol{\ell}) = \psi(\boldsymbol{\varrho})[1 + e^{-i\mathbf{r}\cdot\boldsymbol{\varrho}} - e^{-i\boldsymbol{\ell}\cdot\boldsymbol{\varrho}} - e^{-i(\mathbf{r}-\boldsymbol{\ell})\cdot\boldsymbol{\varrho}}]. \tag{29}$$

Therefore we may write

$$\tilde{\mathcal{A}}^D(\boldsymbol{\varrho},\mathbf{b}) = B\psi(\boldsymbol{\varrho},\alpha) \int \frac{d^2\mathbf{r}}{(2\pi)^2} e^{i\mathbf{r}\cdot\mathbf{b}} \tag{30}$$

$$\times \int \frac{d^2\boldsymbol{\ell}}{\pi\boldsymbol{\ell}^2} \alpha_s(\mu^2)\mathcal{F}(x,x',\boldsymbol{\ell}^2,\mathbf{r}^2)[1 - e^{-i\boldsymbol{\ell}\cdot\boldsymbol{\varrho}}][1 - e^{-i(\mathbf{r}-\boldsymbol{\ell})\cdot\boldsymbol{\varrho}}] \tag{31}$$

$$\equiv \psi(\boldsymbol{\varrho},\alpha) \int \frac{d^2\mathbf{r}}{(2\pi)^2} e^{i\mathbf{r}\cdot\mathbf{b}} \sigma_{q\bar{q}}(\boldsymbol{\varrho},\mathbf{r}), \tag{32}$$

where we have put the $\boldsymbol{\ell}$-dependent part and the factor into the definition of the proton-dipole cross section. If we also carry out the last Fourier transform from momentum transfer to impact parameter, we get the factorised amplitude

$$\tilde{\mathcal{A}}^D(\boldsymbol{\varrho},\mathbf{b}) = \psi(\boldsymbol{\varrho},\alpha)\sigma_{q\bar{q}}(\boldsymbol{\varrho},\mathbf{b}). \tag{33}$$

We see, that also in the case of nonzero momentum transfer we obtain an amplitude, where the photon and proton parts are factorised. Therefore, the the physical picture, where the fluctuation of the virtual photon into the $q\bar{q}$-pair is occuring long before the interaction of dipole and proton, still holds. In turn, the latter depends on the dipole size and now also on the impact parameter $\mathbf{b}$.

4.) In the diffractive scattering of a virtual photon off a proton it is suitable to introduce the photon wave function. Therefore we have calculated the wave functions again, including massive quarks and explicit helicities.

The concept of the photon wave function turned out to arise naturally in considering diffractive scattering, where the factorised part of the photon is expressed in terms of the wave function. Usually the limit of zero momentum transfer is considered here, as in refs. [9, 13]. We saw, that the factorisation of the diffractive amplitude into the photon wave function and the unintegrated structure function still holds, if one includes a momentum transfer.

Finally, we have considered the diffractive amplitude in configuration space. Here, in addition to the dipole size $\boldsymbol{\varrho}$, the dipole cross section will depend on the impact parameter $\mathbf{b}$ of virtual photon and proton in transverse configuration space. We still find consistence with the physical picture, that the fluctuation of the virtual photon into its constituting quarks takes place independently and long before the scattering of the $q\bar{q}$-pair off the proton. This is expressed in the simple, factorised form for the diffractive amplitude in configuration space, in terms of the photon wave function and the dipole cross section.

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